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## Large top mixing from extra dimensions

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### Abstract

Fermion mixing is conveniently described using the effective Lagrangian formalism. We apply this approach to study top mixing in models with an infinite tower of Kaluza-Klein fermion excitations. In the Randall-Sundrum background with a boundary Higgs and phenomenologically viable values of the model parameters, the only effect eventually observable is the loss of universality of the top couplings. Their deviation from the SM predictions can be up to 4% if the five-dimensional Yukawa couplings times the  $\text{AdS}_5$  curvature scale are  $\leq 10$ .

## 1 Introduction

Fermion mixing has been historically a window to new physics and a guide for model building. A primer example was the prediction of the existence of the charm quark to explain the suppression of strangeness-changing neutral currents (GIM mechanism) [1]. Similarly the naturalness of such a suppression for all Flavour-Changing Neutral Currents (FCNC) implies that the isospin properties of the three standard families are the same [2]. However, this may be not what is required to describe the top quark. Its mixing is poorly known experimentally, leaving enough room for relatively large deviations from the Standard Model (SM) predictions (see Ref. [3] for a review). As a matter of fact, there are simple SM extensions where FCNC not involving the top quark are naturally suppressed in a large region of parameter space, and which allow for a large top mixing at the same time [4]. In the following, we first discuss the limits on top mixing [3] and the description of fermion mixing using the effective Lagrangian approach [5]. Then we apply the general results for heavy exotic fermions [6] to top mixing in the five-dimensional Randall-Sundrum (RS) model [7, 8, 9, 10, 11, 12] and show that the size of possible SM deviations can be up to 4%.

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## 2 Top quark couplings

The lowest dimension gauge interactions involving the top quark can be written in standard notation [3, 13, 14]

$$\begin{aligned}
\mathcal{L}_4^{Vtq} = & -g_s \bar{t} \gamma^\mu T^a t G_{\mu a} - \frac{2}{3} e \bar{t} \gamma^\mu t A_\mu \\
& - \frac{g}{\sqrt{2}} \sum_{q=d,s,b} \bar{t} \gamma^\mu (v_{tq}^W - a_{tq}^W \gamma_5) q W_\mu^+ + \text{h.c.} \\
& - \frac{g}{2 \cos \theta_W} \bar{t} \gamma^\mu (v_{tt}^Z - a_{tt}^Z \gamma_5) t Z_\mu \\
& - \frac{g}{2 \cos \theta_W} \sum_{q=u,c} \bar{t} \gamma^\mu (v_{tq}^Z - a_{tq}^Z \gamma_5) q Z_\mu + \text{h.c.}
\end{aligned} \tag{1}$$

The first two terms are fixed by the unbroken gauge symmetry  $SU(3)_C \times U(1)_Q$ , while in the SM the charged currents are parametrized by  $v_{tq}^W, a_{tq}^W = \frac{V_{tq}}{2}$ , where  $V_{tq}$  is the Cabibbo-Kobayashi-Maskawa (CKM) matrix [15], and the neutral currents by  $v_{tt}^Z = \frac{1}{2} - \frac{4}{3} \sin^2 \theta_W$ ,  $a_{tt}^Z = \frac{1}{2}$  and  $v_{tq}^Z, a_{tq}^Z = 0$ . The SM Yukawa couplings are also diagonal. For discussing possible SM extensions using the effective Lagrangian formalism above the electroweak scale it is more convenient, however, to work with left (LH) and right (RH)-handed fields. Thus

$$\begin{aligned}
\mathcal{L}_4^{Vqq'} = & -\frac{g}{\sqrt{2}} (\bar{u}_L^i W_{ij}^L \gamma^\mu d_L^j + \bar{u}_R^i W_{ij}^R \gamma^\mu d_R^j) W_\mu^+ + \text{h.c.} \\
& - \frac{g}{2 \cos \theta_W} \left( \bar{u}_L^i X_{ij}^{uL} \gamma^\mu u_L^j + \bar{u}_R^i X_{ij}^{uR} \gamma^\mu u_R^j \right. \\
& \left. - \bar{d}_L^i X_{ij}^{dL} \gamma^\mu d_L^j - \bar{d}_R^i X_{ij}^{dR} \gamma^\mu d_R^j - 2 \sin^2 \theta_W J_{\text{EM}}^\mu \right) Z_\mu,
\end{aligned} \tag{2}$$

where the summation on family indices  $i, j = 1, 2, 3$  is understood and

$$\begin{aligned}
v^W = \frac{W^L + W^R}{2}, \quad a^W = \frac{W^L - W^R}{2}, \\
v^Z = \frac{X^L + X^R}{2} - 2 \sin^2 \theta_W Q, \quad a^Z = \frac{X^L - X^R}{2},
\end{aligned} \tag{3}$$

with the electric charge  $Q$  term only subtracted for the diagonal couplings.

Present limits on FCNC involving the light quarks are somewhat stringent, but this is not yet the case for the top quark [3, 13, 16]. For instance, from the size of  $K^+ \rightarrow \pi^+ \nu \bar{\nu}$  and neglecting large cancellations with four-fermion contributions  $|X_{ds}^{L,R}| < 4 \times 10^{-5}$ , whereas available collider data only imply  $|X_{tq}^{L,R}| < 0.66$ . Top couplings are constrained by electroweak data within the SM with a higher precision than will be by future direct measurements at large colliders. The question is if one can expect to observe departures from the SM top predictions in future experiments. The answer is positive as we will discuss in last section.

Table 1: Dimension 6 operators contributing to renormalizable  $Vtq$  couplings after SSB.

$$\begin{aligned}
\mathcal{O}_{\phi q}^{(1)} &= (\phi^\dagger i D_\mu \phi) (\bar{q} \gamma^\mu q) & \mathcal{O}_{u\phi} &= (\phi^\dagger \phi) (\bar{q} u \tilde{\phi}) & \mathcal{O}_{\phi W} &= \frac{1}{2} (\phi^\dagger \phi) W_{\mu\nu}^I W^{I\mu\nu} \\
\mathcal{O}_{\phi q}^{(3)} &= (\phi^\dagger \tau^I i D_\mu \phi) (\bar{q} \gamma^\mu \tau^I q) & \mathcal{O}_{d\phi} &= (\phi^\dagger \phi) (\bar{q} d \phi) & \mathcal{O}_{\phi B} &= \frac{1}{2} (\phi^\dagger \phi) B_{\mu\nu} B^{\mu\nu} \\
\mathcal{O}_{\phi u} &= (\phi^\dagger i D_\mu \phi) (\bar{u} \gamma^\mu u) & & & \mathcal{O}_{WB} &= (\phi^\dagger \tau^I \phi) W_{\mu\nu}^I B^{\mu\nu} \\
\mathcal{O}_{\phi d} &= (\phi^\dagger i D_\mu \phi) (\bar{d} \gamma^\mu d) & & & \mathcal{O}_\phi^{(1)} &= (\phi^\dagger \phi) (D_\mu \phi^\dagger D^\mu \phi) \\
\mathcal{O}_{\phi\phi} &= (\phi^T \epsilon i D_\mu \phi) (\bar{u} \gamma^\mu d) & & & \mathcal{O}_\phi^{(3)} &= (\phi^\dagger D^\mu \phi) (D_\mu \phi^\dagger \phi)
\end{aligned}$$

### 3 Top quark effective Lagrangian

The unbroken  $U(1)_Q$  protects the term proportional to  $J_{\text{EM}}^\mu$  in Eq. (2). The other Z and W couplings get contributions from higher orders in the effective Lagrangian expansion after the Spontaneous Symmetry Breaking (SSB) of  $SU(2)_L \times U(1)_Y$ . The lowest order corrections result from the dimension 6 Lagrangian [17]

$$\mathcal{L}_6^{\text{eff}} = \frac{\alpha_x}{\Lambda^2} \mathcal{O}_x + \text{h.c.}, \quad (4)$$

where  $\Lambda$  is the effective high scale and the relevant operators  $\mathcal{O}_x$  are collected in Table 1. As a function of the coefficients  $\alpha_x$  the  $X$  and  $W$  coupling matrices in Eq. (2) read to order  $\frac{v^2}{\Lambda^2}$ , with  $v$  the electroweak vacuum expectation value, [5]

$$\begin{aligned}
X_{ij}^{uL} &= \delta_{ij} - \frac{1}{2} \frac{v^2}{\Lambda^2} V_{ik} (\alpha_{\phi q}^{(1)} + \alpha_{\phi q}^{(1)\dagger} - \alpha_{\phi q}^{(3)} - \alpha_{\phi q}^{(3)\dagger})_{kl} V_{lj}^\dagger, \\
X_{ij}^{uR} &= -\frac{1}{2} \frac{v^2}{\Lambda^2} (\alpha_{\phi u} + \alpha_{\phi u}^\dagger)_{ij}, \\
X_{ij}^{dL} &= \delta_{ij} + \frac{1}{2} \frac{v^2}{\Lambda^2} (\alpha_{\phi q}^{(1)} + \alpha_{\phi q}^{(1)\dagger} + \alpha_{\phi q}^{(3)} + \alpha_{\phi q}^{(3)\dagger})_{ij}, \\
X_{ij}^{dR} &= \frac{1}{2} \frac{v^2}{\Lambda^2} (\alpha_{\phi d} + \alpha_{\phi d}^\dagger)_{ij}, \\
W_{ij}^L &= \tilde{V}_{ik} \left( \delta_{kj} + \frac{v^2}{\Lambda^2} (\alpha_{\phi q}^{(3)})_{kj} \right), \\
W_{ij}^R &= -\frac{1}{2} \frac{v^2}{\Lambda^2} (\alpha_{\phi\phi})_{ij},
\end{aligned} \quad (5)$$

where

$$\tilde{V} = V + \frac{v^2}{\Lambda^2} (V A_L^d - A_L^u V) \quad (6)$$

is a unitary matrix redefining the initial CKM matrix  $V$  to take into account the diagonalization of the quark mass matrices to order  $v^2/\Lambda^2$ .  $A_L^{u,d}$  only involve the

coefficients of the dimension 6 operators  $\mathcal{O}_{u\phi,d\phi}$ , respectively. The non-linear realization of this Lagrangian is studied in Ref. [18].

The couplings in Eq. (5) incorporate features that are forbidden in the SM, namely, FCNC, RH neutral currents not proportional to  $J_{\text{EM}}^\mu$ , RH charged currents, and LH charged currents which are not described by a unitary matrix. These couplings can be directly determined for instance from processes involving the  $V\bar{q}q'$  vertices in which the final  $V$  is observed. For the top quark this will be possible at large colliders [3, 19]. Trilinear couplings also contribute to four-fermion processes (such as mixing of neutral mesons), but four-fermion operators may contribute in this case as well. One can still use these processes to put limits on these couplings assuming that there are no strong cancellations between cubic and quartic couplings.

The operators in the last column of Table 1 redefine the gauge bosons and thus the trilinear couplings in Eq. (5). However, they are flavour blind and hence also constrained by precise electroweak data involving only the light quark flavours, giving then unobservable corrections to top couplings.

## 4 Bulk fermions in the RS model

Models in extra dimensions can give towers of Kaluza–Klein (KK) excitations in four dimensions not far from the TeV scale [20]. The integration of KK gauge bosons and/or fermions generate  $v^2/\Lambda^2$  corrections to the  $X$  and  $W$  couplings in Eq. (5). In both cases they can be originated at tree level and can be *a priori* large [21]. However, the KK gauge boson contributions are proportional to their couplings to the SM fermions times their couplings to the SM gauge bosons, and the latter are typically  $\sim 0.01$  [13, 22]. In this case the corresponding top corrections would be too small to be observable. In the following we will concentrate on the KK fermion contributions, studying in detail the case of bulk fermions in the RS model [11].

In five dimensions there are no chiral fermions. Thus five-dimensional fermions  $\Psi$  are vector-like and can have a Dirac mass of the form

$$\mathcal{L}_D = -im_\Psi(\bar{\Psi}_L\Psi_R + \bar{\Psi}_R\Psi_L), \quad (7)$$

where  $\Psi$  is the sum of the two four-dimensional “chiralities”  $\Psi_{L,R} = \pm\gamma_5\Psi_{L,R}$ , transforming in the same way under the gauge group. The five-dimensional action containing the Yukawa interactions for the three standard families can be written in general with the Higgs on the TeV ( $\pi R$ ) brane

$$\begin{aligned} S_{\text{Yuk}} = & -i \int d^4x \int dy \sqrt{-g} \left[ \lambda_{ij}^{u(5)} \bar{q}_i(x,y) \tilde{u}_j(x,y) \tilde{\phi}(x) \right. \\ & \left. + \lambda_{ij}^{d(5)} \bar{q}_i(x,y) \tilde{d}_j(x,y) \phi(x) + \text{h.c.} \right] \delta(y - \pi R). \end{aligned} \quad (8)$$

Expanding the five-dimensional fields in KK towers and integrating over the fifth dimension, one obtains the four-dimensional mass Lagrangian which after SSB reads

$$\begin{aligned}
i\mathcal{L}_{\text{mass}} = & \sum_{n,m=0}^{\infty} \left[ \lambda_{ij}^{u(nm)} \bar{u}_L^{(n)i} \tilde{u}_R^{(m)j} + \lambda_{ij}^{d(nm)} \bar{d}_L^{(n)i} \tilde{d}_R^{(m)j} \right] + \text{h.c.} \\
& + \sum_{n=1}^{\infty} \left[ M_i^{q(n)} (\bar{u}_L^{(n)i} u_R^{(n)i} + \bar{u}_R^{(n)i} u_L^{(n)i} + \bar{d}_L^{(n)i} d_R^{(n)i} + \bar{d}_R^{(n)i} d_L^{(n)i}) \right. \\
& \quad + M_i^{u(n)} (\bar{\tilde{u}}_L^{(n)i} \tilde{u}_R^{(n)i} + \bar{\tilde{u}}_R^{(n)i} \tilde{u}_L^{(n)i}) \\
& \quad \left. + M_i^{d(n)} (\bar{\tilde{d}}_L^{(n)i} \tilde{d}_R^{(n)i} + \bar{\tilde{d}}_R^{(n)i} \tilde{d}_L^{(n)i}) \right], \tag{9}
\end{aligned}$$

where we have added to Eq. (8) the corresponding Dirac masses in Eq. (7). The latter can always be taken diagonal. The four-dimensional masses can be written

$$\lambda_{ij}^{u,d(nm)} = \lambda_{ij}^{u,d} a_q^{(n)i} a_{u,d}^{(m)j}, \tag{10}$$

with

$$\begin{aligned}
\lambda_{ij}^{u,d} &= \lambda_{ij}^{u,d(5)} k \frac{v}{\sqrt{2}} \sim \text{SM masses}, \\
a_q^{(n)i} &= e^{\pi k R/2} \frac{f_{qL}^{(n)i}(\pi R)}{\sqrt{2\pi k R}}, \\
a_{u,d}^{(m)j} &= e^{\pi k R/2} \frac{f_{u,dR}^{(m)j}(\pi R)}{\sqrt{2\pi k R}},
\end{aligned} \tag{11}$$

and

$$v = e^{-\pi k R} v^{(5)} \sim 250 \text{ GeV}. \tag{12}$$

The factor  $e^{\pi k R}$  in  $\lambda_{ij}^{u,d(nm)}$  is due to the rescaling of the boundary Higgs canonically normalized and the expansion coefficients  $f^{(n)}$  give the  $n$ -th fermion wave functions at the  $\pi R$  brane ( $k \sim 2.44 \times 10^{18} \text{ GeV}$  and  $kR \sim 11$  [12]). It must be noticed that odd fields are zero at the TeV boundary and then the odd chiralities ( $q_R, \tilde{u}_L, \tilde{d}_L$ ) have zero Yukawa couplings for a boundary Higgs. In matrix notation Eq. (9) reads

$$\mathcal{M}^u = \begin{pmatrix} \bar{u}_L^{(0)} \\ \bar{\tilde{u}}_L^{(1)} \\ \vdots \\ \bar{u}_L^{(1)} \\ \vdots \end{pmatrix} \begin{pmatrix} \tilde{u}_R^{(0)} & \tilde{u}_R^{(1)} & \dots & u_R^{(1)} & \dots \\ \lambda_{ij}^u a_q^{(0)i} a_u^{(0)j} & \lambda_{ij}^u a_q^{(0)i} a_u^{(1)j} & \dots & 0 & \dots \\ 0 & M_i^{u(1)} \delta_{ij} & \dots & 0 & \dots \\ \vdots & \vdots & \ddots & \vdots & \vdots \\ \lambda_{ij}^u a_q^{(1)i} a_u^{(0)j} & \lambda_{ij}^u a_q^{(1)i} a_u^{(1)j} & \dots & M_i^{q(1)} \delta_{ij} & \dots \\ \vdots & \vdots & & \vdots & \ddots \end{pmatrix}, \tag{13}$$

and similarly for  $\mathcal{M}^d$ . It is finally convenient in order to compare with experiment, to rotate the zero modes

$$\tilde{u}_R^{(0)i} = (U_R^u)_{ij} \tilde{u}_R^{\prime(0)j}, \quad \tilde{d}_R^{(0)i} = (U_R^d)_{ij} \tilde{d}_R^{\prime(0)j}, \quad q_L^{(0)i} = (U_L^q)_{ij} q_L^{\prime(0)j}, \quad (14)$$

where  $\tilde{u}_R^{\prime(0)}$ ,  $\tilde{d}_R^{\prime(0)}$ ,  $q_L^{\prime(0)}$  are the quark mass eigenstates up to mixing with the heavy KK excitations. In this basis the  $3 \times 3$  light mass submatrices are

$$(U_L^{q\dagger})_{ik} \lambda_{kl}^u a_q^{(0)k} a_u^{(0)l} (U_R^u)_{lj} = V_{ij}^\dagger m_j^u, \quad (U_L^{q\dagger})_{ik} \lambda_{kl}^d a_q^{(0)k} a_d^{(0)l} (U_R^d)_{lj} = m_i^d \delta_{ij}, \quad (15)$$

with  $m_i^{u,d}$  the quark masses and  $V$  the CKM matrix in the absence of mixing. The effective Lagrangian resulting from integrating out the heavy vector-like KK fermions is explicitly given in Refs. [6, 11]. For the RS model the effective couplings in Eq. (5) read

$$\begin{aligned} X_{ij}^{uL} &= \delta_{ij} - m_i^u (U_R^{u\dagger})_{ik} \left[ \sum_{n=1}^{\infty} \left( \frac{a_u^{(n)k}}{a_u^{(0)k}} \right)^2 \frac{1}{M_k^{u(n)/2}} \right] (U_R^u)_{kj} m_j^u, \\ X_{ij}^{uR} &= m_i^u V_{il} (U_L^{q\dagger})_{lk} \left[ \sum_{n=1}^{\infty} \left( \frac{a_q^{(n)k}}{a_q^{(0)k}} \right)^2 \frac{1}{M_k^{q(n)/2}} \right] (U_L^q)_{kr} V_{rj}^\dagger m_j^u, \\ X_{ij}^{dL} &= \delta_{ij} - m_i^d (U_R^{d\dagger})_{ik} \left[ \sum_{n=1}^{\infty} \left( \frac{a_d^{(n)k}}{a_d^{(0)k}} \right)^2 \frac{1}{M_k^{d(n)/2}} \right] (U_R^d)_{kj} m_j^d, \\ X_{ij}^{dR} &= m_i^d (U_L^{q\dagger})_{ik} \left[ \sum_{n=1}^{\infty} \left( \frac{a_q^{(n)k}}{a_q^{(0)k}} \right)^2 \frac{1}{M_k^{q(n)/2}} \right] (U_L^q)_{kj} m_j^d, \\ W_{ij}^L &= \tilde{V}_{ij} - \frac{1}{2} m_i^u (U_R^{u\dagger})_{ik} \left[ \sum_{n=1}^{\infty} \left( \frac{a_u^{(n)k}}{a_u^{(0)k}} \right)^2 \frac{1}{M_k^{u(n)/2}} \right] (U_R^u)_{kl} m_l^u \tilde{V}_{lj} \\ &\quad - \frac{1}{2} \tilde{V}_{il} m_l^d (U_R^{d\dagger})_{lk} \left[ \sum_{n=1}^{\infty} \left( \frac{a_d^{(n)k}}{a_d^{(0)k}} \right)^2 \frac{1}{M_k^{d(n)/2}} \right] (U_R^d)_{kj} m_j^d, \\ W_{ij}^R &= m_i^u V_{il} (U_L^{q\dagger})_{lk} \left[ \sum_{n=1}^{\infty} \left( \frac{a_q^{(n)k}}{a_q^{(0)k}} \right)^2 \frac{1}{M_k^{q(n)/2}} \right] (U_L^q)_{kj} m_j^d, \end{aligned} \quad (16)$$

where at this order  $V$  can be replaced by  $\tilde{V}$  in  $X^{uR}$  and  $W^R$ . The extra contributions are products of  $3 \times 3$  matrices, where the one in square brackets is diagonal and depends on the fermion location and then on the details of the model, and the others are unitary combinations of SM masses. The matrix in the middle can be further simplified noting that

$$(a^{(n)})^2 = (a^{(1)})^2 = 1, \quad (17)$$

which, up to a constant, leaves the diagonal elements as an infinite sum of the inverse of the KK heavy masses  $M^{(n)}$  squared. The lightest mass  $M^{(1)}$  plays the rôle of the effective scale in Eq. (5), and the SM masses include the electroweak vacuum expectation value. The size of the corrections in the RS background depends on one mass parameter per flavour and  $SU(2)_L \times U(1)_Y$  multiplet. One can try to explain the observed pattern of fermion masses in this extended model [23] or as we do in the following, simply ask how large can the SM corrections to the effective couplings be.

## 5 Experimental implications

In popular SM extensions the expected flavour changing branching ratios for the top quark are too small ( $< 10^{-4}$ ) to be observable at future colliders, except in the case of extra vector-like quarks near the electroweak scale (see Table 2 [3]). For SM extensions with exotic quarks [24] electroweak data imply for instance  $|X_{tc}^L| < 0.08$  and  $|X_{tc}^R| < 0.16$  [4], and then branching ratios  $\lesssim 10^{-2}$ . The RS model which is a particular case with an infinite tower of vector-like quark singlets and doublets must satisfy these limits. In fact, the fit of Eq. (16) to present data saturates these bounds. However for  $M^{(1)} \gtrsim 10$  TeV [12] the corresponding Yukawa couplings  $\lambda^{(5)}k$  are  $\sim 100$  and thus too large to recover the successful SM description below the TeV scale. Allowing only for  $\lambda^{(5)}k \leq 10$  the top coupling corrections are reduced, becoming only eventually observable for  $X_{tt}$  and  $W_{tb}^L$ . In this region of parameters  $|\Delta X_{ij}| \sim \frac{m_i m_j}{m_i^2} \Delta X_{tt}^{uL}$  and then the main phenomenological constraints result from the loss of universality of the top couplings [11]. Maximizing the top mixing in this region of parameter space and requiring at least the same  $\chi^2$  as for the SM we have obtained departures

Table 2: Experimental limits expected at LHC for the top quark flavour changing branching ratios  $\text{Br}(t \rightarrow uZ, cZ)$  and values predicted in the SM, the two Higgs model (2H), supersymmetric models (SUSY) without R and with  $\tilde{R}$  parity breaking and the SM extensions with exotic quarks. The branching ratio is defined as the decay rate divided by 1.56 GeV.

LHC	SM	2H	SUSY $\tilde{R}, R$	Vector – like quarks
$\sim 10^{-4}$	$\sim 10^{-13}$	$\sim 10^{-6}$	$\sim 10^{-4}, 10^{-8}$	$\sim 10^{-2}$

from the SM top couplings of up to  $\sim 4\%$ . Indeed for the Yukawa couplings

$$\lambda_{ij}^{u(5)} k = \begin{pmatrix} 5.6 \times 10^{-4} & 10 & 9.2 \times 10^{-3} \\ -6.2 \times 10^{-4} & -10 & -9.2 \times 10^{-3} \\ 1.2 \times 10^{-3} & -10 & 1.8 \times 10^{-2} \end{pmatrix}, \quad (18)$$

$$\lambda_{ij}^{d(5)} k = \begin{pmatrix} -5.9 \times 10^{-5} & -1.1 \times 10^{-3} & 4.2 \times 10^{-2} \\ -7.7 \times 10^{-5} & 5.2 \times 10^{-4} & -4.2 \times 10^{-2} \\ 1.8 \times 10^{-5} & -1.5 \times 10^{-3} & -4.6 \times 10^{-2} \end{pmatrix} \quad (19)$$

and the coefficients

$$\begin{aligned} a_u^{(0)u,c,t} &= 0.5603, 0.1018, 0.5603, \\ a_d^{(0)i} &= a_q^{(0)i} = 0.5603, \end{aligned} \quad (20)$$

where  $i = 1, 2, 3$  is the family index, we find

$$X_{tt}^L = 0.9608, \quad X_{tt}^R = 0.0013, \quad (21)$$

and

$$|W_{ij}^L| = \begin{pmatrix} 0.9752 & 0.2227 & 0.0036 \\ 0.2227 & 0.9752 & 0.0402 \\ 0.0096 & 0.0394 & 0.9804 \end{pmatrix}. \quad (22)$$

The corrections to the remaining  $X$  and  $W$  couplings are negligible. We have taken  $k = 2.44 \times 10^{18}$  and  $kR = 10.815$  [12]. The rotation matrices  $U$  entering in Eq. (16) are fixed by the standard fermion masses and the CKM matrix through Eq. (15). Note that the value  $a^{(0)} = 0.5603$  in Eq. (20) ensures that the bound on the lightest KK gauge boson mass  $M_1^{\text{gauge}} \gtrsim 10$  TeV is satisfied. This is needed to suppress the contribution of the tower of KK gauge bosons to electroweak observables below present experimental limits [10].

In summary, we have shown that in the RS model there can be deviations from the SM  $Z\bar{t}t$  and  $W\bar{t}b$  couplings of at most  $\sim 4\%$  and  $2\%$ , respectively, due to the top mixing with the KK fermion excitations. This has to be compared with the measurement of  $W_{tb}^L$  at LHC, for which an accuracy of  $5\%$  ( $10\%$  for  $|W_{tb}^L|^2$  in the cross section) is foreseen as an ambitious but attainable goal [3]. There are better prospects for the measurement of  $X_{tt}^L$  at TESLA. For instance, with a center of mass energy of 500 GeV, an integrated luminosity of  $300 \text{ fb}^{-1}$  and unpolarized beams one expects to collect in the detector 34800 top pairs with one  $W$  decaying into  $e\nu$  or  $\mu\nu$  and the other  $W$  decaying hadronically, reaching a precision of  $2\%$  in the determination of the  $Z\bar{t}t$  coupling [25]. This precision will improve when all channels are included.

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